Structure of Low-Energy Collective 0^--States in Doubly Magic Nuclei and Matrix Elements of the P-odd and P- and T-odd Weak Interaction

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Abstract

The structure of the collective low-energy $J^\pi = 0^-$ (T=0 and T=1) modes is studied for a doubly magic nucleus in a schematic analytic model of RPA. The 0^- phonon states ($T = 0, 1$) lie at energies $E_{T=0}(0^-) \lesssim \omega$ and $E_{T=1}(0^-) > \omega$, where ω is the oscillator frequency. The matrix elements of P-odd and P- and T-odd weak one-body potentials connecting the ground state to these 0^--states, W_{coll}, are enhanced by the factor $\sim 2(\frac{\omega}{E})^{1/2} A^{1/3} \sim 10$ as compared to the single-particle value w_{sp} what can result in values $|W_{coll}| \sim 20 - 30eV$ if standard values of DDH parameters are used for w_{sp}. Similar enhancement arises in the P- and T-odd case.

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The purpose of this work is to study low-lying collective nuclear states with quantum numbers $J^\pi = 0^-$ and the matrix elements of parity nonconserving potentials connecting these states to the ground state. The work is motivated not only by the interest in the low-energy 0^- modes, which already attracted attention earlier in relation to pion (pre-)condensation, but also by the possibility of the collective enhancement of P- and P-,T-odd matrix elements of the pseudoscalar "weak potentials" in nuclei. The importance of collective 0^- states for these effects has been already revealed within doorway state approach to the P-violation in compound nuclei [1], [2]. It was pointed out [1] that the 0^- component of the spin-dipole state is an important mediator of parity mixing between nuclear states. This applies to both parity violating (P-odd) and time reversal conserving (T-even) as well as to P-odd and time reversal violating (T-odd) interactions. The one-body potentials for these two symmetry violating interactions can be approximately described by the operators $(\vec{\sigma}\vec{p})$ for P-odd T-even and $(\vec{\sigma}\vec{r})$ for P-odd, T-odd. When acting on a nuclear state both operators will excite the 0^- mode. Since the weak interaction does not conserve isospin one considers both isoscalar $(T = 0)$ and $(T = 1)$ $J = 0^-$ spin-dipole states.

In Refs. [4,5], renormalization of the weak P- and P-,T-odd potentials due to the residual strong interaction were calculated and the positions of low energy 0^- modes was discussed. The properties of 0^- $(T = 1)$ states were studied earlier in an elaborated Hartree-Fock+RPA approach in Ref. [3]. Spin-dipole modes were studied in [6], [8] within an approximate RPA approach. An intersting analysis of related effects is presented in Ref. [9].

Here we show, using an approximate analytical RPA calculation, that a realistic residual nucleon-nucleon interaction which has been widely used to describe various nuclear properties does indeed lead to formation of such collective 0^- modes. We give approximate quantitative predictions for their characteristics. We confine ourselve here to the case of doubly magic nucleus. The density of states is relatively low. We neglect in our study the coupling of the one-particle one-hole 0^- states to all other configurations.

We write the nuclear Hamiltonian in the form

$$H = H_0 + V, \tag{1}$$

where V is the residual strong interaction and H_0 is the singe-particle shell-model Hamiltonian of nucleons with momenta $\vec{p}_a$ and mass m moving in a spherically symmetric average potential $U(r)$:

$$H_0 = \sum_a \left(\frac{\vec{p}_a^2}{2m} + U(r_a) \right), \qquad U(r) = -|U_0| + \frac{m\omega^2 \vec{r}_a^2}{2} \tag{2}$$

(we use the harmonic oscillator well of the depth $|U_0|$ and neglect for simplisity the shell splitting).

We employ the Landau-Migdal parametrization [10] of the residual strong interaction chosing V to be the sum of velocity independent V_0 and velocity dependent V_1 components, $V = V_0 + V_1$. The velocity independent term in the residual interaction is given by

$$V_0(1,2) = C\delta(\vec{r}_1 - \vec{r}_2)[f_0 + f_0'\vec{\tau}_1\vec{\tau}_2 + g\vec{\sigma}_1\vec{\sigma}_2 + g'\vec{\tau}_1\vec{\tau}_2\vec{\sigma}_1\vec{\sigma}_2], \tag{3}$$

where $C = 300\ MeV \cdot fm^3$ is the universal Migdal constant. (We use the convention $\tau_z(p) = -1, \tau_z(n) = 1$ [11]).

The spin- and momentum-dependent component of the interaction V_1 will be used in two forms: the first one can be written [12] in the form of a Landau-Migdal spin- and velocity dependent interaction,

$$V_1(1,2) = \frac{1}{4}Cp_F^{-2}(g_1 + g_1'\vec{\tau}_1\vec{\tau}_2)(\vec{\sigma}_1\vec{\sigma}_2)[\vec{p}_1\vec{p}_2\delta(\vec{r}_1 - \vec{r}_2) + \vec{p}_1\delta(\vec{r}_1 - \vec{r}_2)\vec{p}_2 + \vec{p}_2\delta(\vec{r}_1 - \vec{r}_2)\vec{p}_1 + \delta(\vec{r}_1 - \vec{r}_2)\vec{p}_1\vec{p}_2], \tag{4}$$

p_F is the Fermi momentum. In Eqs.(3,4), the parameters f and g are the strength constants of order of unity (their values will be given below). Generally speaking, they contain r-dependence which is usually parametrized via the density dependence. However, this dependence is considerable only for the constants f and f' which describe spin-independent components of the interaction. In what follows, only spin-dependent components of the interaction are essential, and the corresponding strengths g_i, g_i' are density independent. This

phenomenological interaction can be viewed as a model of more realistic but also more complicated forces. It is therefore convenient to use it first because of simplicity of calculations. In fact, the results obtained for this interaction manifest the correct tendencies (see also [4]).

We also make calculations for the π- and ρ- meson exchange interaction in its explicit form

$$V_{\pi+\rho} = \frac{f_\pi^2}{m_\pi^2}(\vec{\tau}_1\vec{\tau}_2)\Big[(\vec{\nabla}\cdot\vec{\sigma}_1)(\vec{\nabla}\cdot\vec{\sigma}_2), \frac{e^{-m_\pi r}}{r}\Big] - \frac{f_\rho^2}{m_\rho^2}(\vec{\tau}_1\vec{\tau}_2)\Big[(\vec{\nabla}\times\vec{\sigma}_1)(\vec{\nabla}\times\vec{\sigma}_2), \frac{e^{-m_\rho r}}{r}\Big] \tag{5}$$

using it instead of V_1 in Eq.(4). Here, m_π (m_ρ) are the π-meson (ρ-meson) mass, f_π and f_ρ are the coupling constants, and $\vec{\nabla}$ is the derivative with respect to the nucleon separation $\vec{r} = \vec{r}_1 - \vec{r}_2$; $\times$ denotes the external vector product. The standard “Lorentz-Lorentz” correction (see, e.g., [13]) is assumed [4].

In the general case, the RPA equations can be written in the form of equations of motion for the phonon creation operators $\hat{A}^\dagger$

$$-E_n\hat{A}_n^\dagger = [\hat{A}_n^\dagger, \mathcal{H}_{RPA}] = [\hat{A}_n^\dagger, H_0] + \langle[\hat{A}_n^\dagger, V]\rangle, \tag{6}$$

where the first equality comes from the commutator of the phonon creation operator $\hat{A}_n^\dagger$ with the RPA Hamiltonian $\mathcal{H}_{RPA} \equiv const + \sum_n E_n\hat{A}_n^\dagger\hat{A}_n$ which is derived from the initial Hamiltonian $H_0 + V$ when the RPA phonons are used. The expectation value in (6) is evaluated using the uncorrelated ground state $|0^+\rangle$. Within the RPA, the phonon operators are

$$\hat{A}_n^\dagger = \sum_{im} X_{mi}^n \hat{c}_m^\dagger \hat{c}_i - \sum_{im} Y_{mi}^n \hat{c}_i^\dagger \hat{c}_m \tag{7}$$

where the $\hat{c}_a^\dagger$ and $\hat{c}_b$ are the operators of the creation and annihilation of nucleons in the single-particle states marked by a and b, respectively. The amplitudes X and Y are related to the “forward” and “backward” going graphs, respectively. The index n marks the modes, the total number of which $\mathcal{N}$, coincides with the total number of the particle-hole pairs with quantum numbers 0^-.

The 0^- particle-hole states have energies $1\omega, 3\omega, ...$. We are interested in low-energy collective states which should be dominated by the 1ω transitions. The corresponding constituents of the phonon modes can be conveniently parametrized by the elementary correlated 0^- particle-hole operators $\hat{A}_0^\dagger$ and $\hat{A}_0$ defined below. We use the following representation for the *collective* phonon creation and destruction operators $\hat{A}_{coll}^\dagger$ and $\hat{A}_{coll}$

$$\hat{A}_{coll}^\dagger = x\sum_{mi}(\vec{\sigma}\vec{a}^\dagger)_{mi}\hat{c}_m^\dagger\hat{c}_i - y\sum_{mi}(\vec{\sigma}\vec{a})_{im}\hat{c}_i^\dagger\hat{c}_m, \tag{8}$$

and $\hat{A}_{coll}$ given by the Hermitean conjugate of $\hat{A}_{coll}^\dagger$. Here, $(O)_{ab}$ means the matrix elements of a single-particle operator O between the nucleon states ψ_a and ψ_b. The symbols $\vec{a}^\dagger$ and $\vec{a}$ denote the harmonic oscillator "raising" and "lowering" operators $\vec{a}^\dagger = \frac{1}{\sqrt{2}}\left[(m\omega)^{1/2}\vec{r} - i(m\omega)^{-1/2}\vec{p}\right]$, $\vec{a} = \frac{1}{\sqrt{2}}\left[(m\omega)^{1/2}\vec{r} + i(m\omega)^{-1/2}\vec{p}\right]$. It is seen that the role of the operators

$$\hat{A}_0^\dagger \equiv \sum_{mi}(\vec{\sigma}\vec{a}^\dagger)_{mi}\hat{c}_m^\dagger\hat{c}_i, \qquad \hat{A}_0 \equiv \sum_{mi}(\vec{\sigma}\vec{a})_{im},\hat{c}_i^\dagger\hat{c}_m, \tag{9}$$

is to create and destroy the correlated particle-hole excitations with the total quantum numbers $J^\pi = 0^-$. Indeed, the operator $a^\dagger$ raises the oscillator principal quantum number N by unity, while the role of the spin operator in the internal product is to rotate properly the nucleon spins. The combined effect is to remove particles from the orbital with the quantum numbers $[N, l, j]$ and place them in the orbital $[N+1, l\pm 1, j]$. Thus, due to this construction the operator $\hat{A}_0^\dagger$ can only *create* the particle-hole pairs of 0^- type (m and i denote the states above and below the Fermi level, correspondingly). The effect of the *annihilation* operator $\hat{A}_0$ is just the opposite to that of $\hat{A}_0^\dagger$.

The expressions for the phonon creation and destruction operators $\hat{A}^\dagger$ and $\hat{A}$ are an ansatz for the more general expression of collective RPA phonon operators given in Eq.(7). The form of the ansatz is based on the simple commutator properties for the operators $(\vec{\sigma}\vec{r})$ and $(\vec{\sigma}\vec{p})$ [4]. with the interactions V_0 and V_1. (Note that if we substitute the interactions V_0 and V_1 by factorized expressions $V_0 \to const(\vec{\sigma}\vec{r})_{ab}(\vec{\sigma}\vec{r})_{cd}$, and $V_0 \to const(\vec{\sigma}\vec{p})_{ab}(\vec{\sigma}\vec{p})_{cd}$, the ansatz (8) will be an exact solution of the RPA equations (6)).

The use of Eq.(8) allows us to get a closed system of RPA equations for the decoupled collective phonon modes ($T = 0, 1$), and instead of system of a large number of coupled equations (6) we obtain the simplified, approximate sytem of the equations for the amplitudes x and y

$$-E_{coll}\hat{A}^{\dagger}_{coll} \simeq [\hat{A}^{\dagger}_{coll}, H_0] + \langle[\hat{A}^{\dagger}_{coll}, V_0 + V_1]\rangle, \qquad -E_{sp}\hat{A}^{\dagger}_{sp} \simeq [\hat{A}^{\dagger}_{sp}, H_0] \tag{10}$$

From the first one of Eqs.(10), the closed equations for the collective phonon amplitudes x and y can be easily obtained using the following relations

$$\langle[\hat{A}^{\dagger}_0, \hat{\mathcal{V}}_0]\rangle = -\frac{\rho_0}{|U_0|}\frac{\omega}{2}(\hat{A}_0 + \hat{A}^{\dagger}_0), \qquad \langle[\hat{A}^{\dagger}_0, \hat{\mathcal{V}}_1]\rangle = -2m\rho\omega(\hat{A}^{\dagger}_0 - \hat{A}_0). \tag{11}$$

where $\hat{\mathcal{V}}_0$ and $\hat{\mathcal{V}}_1$ denote the interaction terms $(\vec{\sigma}_1\vec{\sigma}_2)\delta(\vec{r}_1 - \vec{r}_2)$ and $(\vec{\sigma}_1\vec{\sigma}_2)\{\vec{p}_1, \{\vec{p}_2, \delta(\vec{r}_1 - \vec{r}_2)\}\}$ respectively expressed in the second quantization form (the curly brackets denote an anti-commutator). Hermitean conjugation gives the corresponding relations for the commutators of $\hat{A}_0$ operators. In the first one of Eqs.(11) we assumed the proportionality of the nucleon density ρ and the potential U: $\rho \approx -\frac{\rho_0}{|U_0|}U$, where ρ_0 and U_0 are the values of the density and the potential in the center of the nucleus.

All the remaining phonon degrees of freedom, marked by the subscript 'sp' in Eqs.(10), are assumed here to be essentially noncollective. The contribution from the interaction term in the second one of Eqs.(10) is dropped. The frequencies of those predominantly single-particle modes are not shifted from the shell-model value ω.

For the case of two kinds of nucleons (p and n), the collective phonons operators are sought in the form $\hat{A}^{\dagger}_{coll} = x_p\hat{A}^{\dagger}_0(p) + x_n\hat{A}^{\dagger}_0(n) - y_p\hat{A}_0(p) - y_n\hat{A}_0(n)$ instead of (8). From (10), we obtain the following system of equations for the amplitudes x_p, x_n, y_p and y_n

$$\begin{pmatrix} E_T & 0 & -\omega(1 + \frac{C\rho Z}{|U|A}g^{(+)}) & -\omega\frac{C\rho N}{|U|A}g^{(-)} \\ 0 & E_T & -\omega\frac{C\rho Z}{|U|A}g^{(-)} & -\omega(1 + \frac{C\rho N}{|U|A}g^{(+)}) \\ \omega(1 + \frac{C\rho m Z}{p_F^2 A}g_1^{(+)}) & \omega\frac{C\rho m N}{p_F^2 A}g_1^{(-)} & -E_T & 0 \\ \omega\frac{C\rho m Z}{p_F^2 A}g_1^{(-)} & \omega\frac{C\rho m N}{p_F^2 A}g_1^{(+)} & 0 & -E_T \end{pmatrix} \begin{pmatrix} x_p - y_p \\ x_n - y_n \\ x_p + y_p \\ x_n + y_n \end{pmatrix} = 0, \tag{12}$$

where $g^{(\pm)} = g \pm g'$ and $g_1^{(\pm)} = g_1 \pm g'_1$.

It is convenient to analyze in detail the case of equal numbers of protons and neutrons ($N = Z$). From the consistency condition for (12) we find the two eigenfrequencies of the collective modes $E_{T=0}$ and $E_{T=1}$ to be

$$E_{T=0} = \omega\Big[\Big(1 + \frac{C\rho m}{p_F^2} g_1\Big)\Big(1 + \frac{C\rho}{|U|} g\Big)\Big]^{1/2}, \quad E_{T=1} = \omega\Big[\Big(1 + \frac{C\rho m}{p_F^2} g_1'\Big)\Big(1 + \frac{C\rho}{|U|} g'\Big)\Big]^{1/2}, \tag{13}$$

The eigenvectors of the system (12) are easily found using Eq.(13), in a form of phonon creation operators

$$\begin{aligned}
\hat{A}^\dagger_{T=0} = \frac{1}{2}\Big[\frac{\omega}{E_{T=0}(1 + \frac{C\rho m}{p_F^2} g_1)\mathcal{N}_{coll}}\Big]^{1/2} \Big\{ \Big[\frac{E_{T=0}}{\omega} + 1 + \frac{C\rho m}{p_F^2} g_1\Big](\hat{A}^\dagger_0(p) + \hat{A}^\dagger_0(n)) \\
- \Big[-\frac{E_{T=0}}{\omega} + 1 + \frac{C\rho m}{p_F^2} g_1\Big](\hat{A}_0(p) + \hat{A}_0(n))\Big\}, \\
\hat{A}^\dagger_{T=1} = -\frac{1}{2}\Big[\frac{\omega}{E_{T=1}(1 + \frac{C\rho m}{p_F^2} g_1')\mathcal{N}_{coll}}\Big]^{1/2} \Big\{ \Big[\frac{E_{T=1}}{\omega} + 1 + \frac{C\rho m}{p_F^2} g_1'\Big](\hat{A}^\dagger_0(p) - \hat{A}^\dagger_0(n)) \\
+ \Big[-\frac{E_{T=1}}{\omega} + 1 + \frac{C\rho m}{p_F^2} g_1'\Big](\hat{A}_0(p) - \hat{A}_0(n))\Big\},
\end{aligned} \tag{14}$$

where $\mathcal{N}_{coll}$ is the colective phonon normalization constant. The operators of the independent RPA modes should obey the orthogonality and normalization conditions, thus

$$[\hat{A}^\dagger_T, \hat{A}^\dagger_{T'}] = 0, \quad \langle 0^+ [\hat{A}_T, \hat{A}^\dagger_{T'}] 0^+\rangle = \delta_{T,T'}. \tag{15}$$

where the expectation value is taken over the uncorrelated ground state. The first set of Eqs.(15) (orthogonality) holds by construction, while the second one determines the normalization. Substituting (14) into (15) gives us the normalization relation

$$\langle 0^+ | [\hat{A}_T, \hat{A}^\dagger_{T'}] | 0^+\rangle = \delta_{T,T'}(R_p + R_n)\mathcal{N}^{-1}_{coll}.$$

Here, for a given kind of nucleons (p or n), the quantity R is

$$R = \langle [(\hat{\vec{\sigma}\vec{a}}), (\hat{\vec{\sigma}\vec{a}^\dagger})]\rangle = \langle \sum_i (3 + 2(\vec{\sigma}\vec{l})_{ii})\hat{c}^\dagger_i \hat{c}_i\rangle = 3\mathcal{N}$$

where $\vec{l}$ is the single-particle orbital angular momentum. The expectation value here is reduced to the summation over the occupied states i and gives three times the number of particles of given kind, $\mathcal{N}$. Note that for the spin saturated shell-model ground state of a

doubly magic nucleus, the term proportional to $(\vec{\sigma}\vec{l})_{ii}$ does not contribute. Therefore, the phonon normalization constant (for $N = Z = A/2$) is

$$\mathcal{N}_{coll} = 3A. \tag{16}$$

The matrix elements of the operators $(\vec{\sigma}\vec{r})_{p,n}$ and $(\vec{\sigma}\vec{p})_{p,n}$ between the correlated ground state $|0^+)$ and the excited single-phonon states $|0^-, T = 0, 1) = \hat{A}^\dagger_{T=0,1}|0^+)$ can be easily obtained by expressing the operators in terms of the phonon creation and destruction operators $\hat{A}^\dagger_{T=0,1}$ and $\hat{A}_{T=0,1}$ [see Eqs. (9),(14)]. The results for the matrix elements of the isoscalar $W^P_0 = (0^+|(\vec{\sigma}\vec{p})\hat{1}|0^-, T = 0), W^{PT}_0 = (0^+|(\vec{\sigma}\vec{r})\hat{1}|0^-, T = 0)$ and isovector parity violating operators $W^P_1 = (0^+|(\vec{\sigma}\vec{p})\hat{1}|0^-, T = 1)$, $W^{PT}_1 = (0^+|(\vec{\sigma}\vec{r})\hat{1}|0^-, T = 1)$ are found to be:

$$W^{PT}_0 = \Big(\frac{(1 + \frac{Cm\rho}{p_F^2}g_1)R}{2mE_{T=0}}\Big)^{1/2}, \qquad W^P_0 = -i\Big(\frac{mE_{T=0}R}{2(1 + \frac{Cm\rho}{p_F^2}g_1)}\Big)^{1/2},$$
$$W^{PT}_1 = \Big(\frac{(1 + \frac{Cm\rho}{p_F^2}{g'}_1)R}{2mE_{T=1}}\Big)^{1/2}, \quad W^P_1 = -i\Big(\frac{mE_{T=1}R}{2(1 + \frac{Cm\rho}{p_F^2}g'_1)}\Big)^{1/2}. \tag{17}$$

To reveal the collective enhancement factors, we compare these results to the matrix elements of the corresponding operators between shell-model states. For the $(\vec{\sigma}\vec{p})$ and a given 1particle-1hole state, $\hat{c}^\dagger_m\hat{c}_i|0^+\rangle$, we have the following relations

$$\langle 0^+|\widehat{(\vec{\sigma}\vec{p})} \quad \hat{c}^\dagger_m\hat{c}_i|0^+\rangle = -i\left(\frac{m\omega}{2}\right)^{1/2}(\vec{\sigma}\vec{a})_{mi} \quad \sim \quad \left(\frac{m\omega}{2}\right)^{1/2}\sqrt{N_F} \sim \left(\frac{m\omega}{2}\right)^{1/2}(3A)^{1/6} \tag{18}$$

where the last expressions indicate the order of magnitude. N_F denotes the principal oscillator quantum number at the Fermi energy of a doubly magic nucleus. We used in (18) the fact that $N_F \sim (3A)^{1/3}$. Similarly:

$$\langle 0^+|\widehat{(\vec{\sigma}\vec{r})} \quad \hat{c}^\dagger_m\hat{c}_i|0^+\rangle = \left(\frac{1}{2m\omega}\right)^{1/2}(\vec{\sigma}\vec{a})_{mi} \quad \sim \left(\frac{1}{2m\omega}\right)^{1/2}\sqrt{N_F} \sim \left(\frac{1}{2m\omega}\right)^{1/2}(3A)^{1/6} \tag{19}$$

Comparing the matrix elements (17) to the single-particle ones (18), (19) and taking into account Eq.((13), we see that the matrix elements of the P-odd potentials connecting the correlated ground state to the single-phonon states are enhanced by the factors

$$F^{P}_{coll} \approx \frac{|(0^+|(\vec{\sigma}\vec{p})|0^-, T=0)|}{|\widehat{(\vec{\sigma}\vec{p})}_{mi}|} \sim \quad \sqrt{\frac{\omega}{E_0}(1+\frac{C\rho}{|U|}g)} \quad (3A)^{1/3}. \tag{20}$$

and

$$F^{PT}_{coll} \approx \frac{|(0^+|(\vec{\sigma}\vec{r})|0^-, T=0)|}{|\widehat{(\vec{\sigma}\vec{r})}_{mi}|} \sim \quad \sqrt{\frac{\omega}{E_0}(1+\frac{Cm\rho}{p_F^2}g_1)} \quad (3A)^{1/3}. \tag{21}$$

Analogous expressions can be written for the $T=1$ matrix elements by substituting $g'(g'_1)$ instead of $g(g_1)$ in (20,21).

The total collective enhancement factor F_{coll} has a natural form typical for the low-energy collective modes. On the right hand side of Eqs.(20),(21), the second factor is the collective "phase volume" enhancement, F_{pv}, which is roughly equal to the square root of the number of single-particle constituents of the phonon, i.e., approximately the number of nucleons in the upper major shell, $\mathcal{A}_F$),

$$F_{pv} \sim \sqrt{\mathcal{A}_F} \sim N_F \sim (3A)^{1/3}.$$

The first factor in Eqs.(20),(21) is the "adiabaticity parameter" [14]. It reflects the change of the frequency of the collective motion compared to the frequency of single-particle motion ω

$$F_{ad} \sim \sqrt{\frac{\omega}{E_0}(1+\frac{C\rho}{|U|}g)} \qquad (P-odd), \qquad F_{ad} \sim \sqrt{\frac{\omega}{E_0}(1+\frac{Cm\rho}{p_F^2}g_1)} \qquad (P-,T-odd).$$

Note that for the low-lying 0^--mode ($T=0$), $F_{ad} > 1$ gives additional enhancement.

We turn now to a quantitative discussion of the above results. Using the conventional constants of the Landau-Migdal interaction (3),(4) $g=0.575, g'=0.725$, $g_1=-0.5$ and $g'_1=-0.26$ (see e.g. Refs. [10], [15], [16], [12]), we find from Eq.(13):

$$E_{T=0} = 1.03\omega, \qquad E_{T=1} = 1.20\omega. \tag{22}$$

The dependence of E_T on the total strength of the interaction constant is shown on Fig.1. It is seen that due to the negative value of g_1, g'_1, the effect of the velocity dependent interaction is to reduce the frequency of the isoscalar mode pushing it towards instability

[5]. However, the velocity independent component acts in the opposite direction. The excitation energy of the isovector mode that is governed by the primed constants tends to increase with the interaction strength. (These tendencies are in agreement with those found in numerical studies [3] and with results for the meson exchange interactions [4]). The combined and opposite action of V_0 and V_1 on E_T tends to moderate energy shift from ω, so the values (22) do not mean that the correlations in the ground state are small. The collective enhancement of the matrix elements (17) are $F^P \sim 10-15$ for the above values of g.

We should stress that it is not the prime purpose of this work to achieve agreement with experiment. The mean field that provides the basis of the single-particle energies is not treated in this work. Therefore effects of consistency between the $p-h$ interaction and the mean field [15] are not taken into account. Also, we do not include the one-body spin-orbit interaction which for 0^- mode could be important especially in lighter nuclei [6].

As we mentioned, the velocity dependent Landau-Migdal interaction can be viewed as a model of a more realistic momentum-dependent π- and ρ-meson exchange interactions. Using $V_{\pi+\rho}$ of Eq.(5) as a residual interaction within the above formalism we find

$$E_{T=0} = \omega\Big[\Big(1-\frac{3k}{2}\Big)\Big(1+\frac{C\rho}{|U|}g-3\xi\Big)\Big]^{1/2},$$
$$E_{T=1} = \omega\Big[\Big(1+\frac{k}{2}\Big)\Big(1+\frac{C\rho}{|U|}g'-2\phi+\xi\Big)\Big]^{1/2}, \qquad (23)$$

Here, the constants k and ϕ and ξ are related to the meson-nucleon coupling strenghts: $k = 2\rho_0 m q$ results from the evaluation of the exchange terms [4] with $q = 6\pi(\frac{f_\pi^2}{m_\pi^4}W_\pi - \frac{4}{3}\frac{f_\rho^2}{m_\rho^4}W_\rho)$ and the nonlocality factors W ($W_{\pi,\rho} \to 1$ for $m_{\pi,\rho} \to \infty$) are $W_\pi = 0.11$, $W_\rho = 0.69$ for the pion and ρ-meson correspondingly [4]. ϕ results from the direct terms, $\phi = \frac{v}{m\omega^2}$, $v = -\frac{4\pi}{3}\frac{\rho_0}{|U_0|}m\omega^2(\frac{f_\pi^2}{m_\pi^2} + 2\frac{f_\rho^2}{m_\rho^2})$, while the term ξ is due to the exchange terms and $\xi = \frac{2\pi}{3}\frac{\rho_0}{|U_0|}(-\frac{f_\pi^2}{m_\pi^2}{W'}_\pi + 2\frac{f_\rho^2}{m_\rho^2}{W'}_\rho)$ where ${W'}_\pi = 0.56$ and ${W'}_\rho = 0.71$. The pion constant is equal $f_\pi^2 = 0.08$, while for the ρ-meson the coupling constant is in range from $f_\rho^2 = 1.86$ ("weak ρ-meson") to $f_\rho^2 = 4.86$ ("strong ρ-meson") [13], [17].

The phonon operators $\hat{A}_T^\dagger$, $\hat{A}_T$ are given now by the same expressions as before [Eq.(14)]

but with the replacements $1+\frac{C\rho m}{p_F^2}g_1 \to 1-3k/2$ and $1+\frac{C\rho m}{p_F^2}g_1' \to 1+k/2$ and with $E_{T=0,1}$ given by (23). Correspondingly, the matrix elements of the symmetry violating operators

$$
\begin{aligned}
W_0^{PT} &= \Big(\frac{(1-3k/2)R}{2mE_{T=0}}\Big)^{1/2}, \quad W_0^P = -i\Big(\frac{mE_{T=0}R}{2(1-3k/2)}\Big)^{1/2}, \\
W_1^{PT} &= \Big(\frac{(1+k/2)R}{2mE_{T=1}}\Big)^{1/2}, \quad W_1^P = -i\Big(\frac{mE_{T=1}R}{2(1+k/2)}\Big)^{1/2}.
\end{aligned} \tag{24}
$$

The results for the frequencies are plotted in Fig.1. It is seen that the effect of the meson exchange interaction is to push down the isoscalar frequency $E_{T=0}$, mainly due to the π-meson contribution. At the same time it shift upwards the isovector frequency. Similar behaviour of the $E_{T=1}$ was obtained in a more realistic RPA calculations in Ref. [3]. Fig.2a shows E_T as a function of f_ρ^2. Instability of the $0^-, T=0$ occurs for the "weak ρ-meson", as was discussed in [4] using general properties of the linear response [18]. It is interesting that the collapse occurs also for the "strong" ρ-meson, while the stable energy range $E_0 \simeq (0.2-0.4)\omega$ occurs for $1.86 < f_\rho^2 < 4.86$. Fig. 2b presents the enhancement factors F^P and F^{PT} of the weak matrix elements modified according to (24). The collective enhancement for the P-odd matrix elements rises from $10-20$ to considerably larger values as $E_{T=0}^2 \to 0$.

The enhancement of the P-odd matrix elements connecting the ground state to the single-phonon states found here can result in observable P-effects [e.g., in the electromagnetic transitions (multipole mixing) [20]]. For the P-odd ($T=0$) matrix element, where the enhancement is maximum, the quantity W_{coll}^P can reach, for example for ^{206}Pb, a value of

$$
W_{coll}^P = \quad (0^+|g_p^W \frac{G}{2m\sqrt{2}}\{(\vec{\sigma}_p\vec{p}_p),\rho\} + g_n^W \frac{G}{2m\sqrt{2}}\{(\vec{\sigma}_n\vec{p}_n),\rho\}|0^-, T=0) \quad \simeq \quad 20-30eV,
$$

for the realistic weak interaction strengths $g_p \simeq 4$ and $g_n \simeq 1$ (see, e.g., [19]) consistent with the standard parameters of the two-body weak interaction [21], [20] (G is the Fermi constant).

The collective of 0^- states has been already considered in the doorway mechanism approach [1], [2] in the problem of P-odd mixing in compound nuclei. The enhancement in the parity violating matrix element due to the collectivity of the 0^- was one of the assumptions

in the above works. The results of the present work confirm their idea and the collectivity of the P-odd matrix element is a natural result of the RPA treatment of the 0^- states.

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Figure captions

Fig. 1. The 0^- energies $E_{T=0}$ and $E_{T=1}$ [Eq.(13), in units of ω] as a function of the strength of the Landau-Migdal interaction C [taken in units of the standard value (Eq.3,4)] (solid curves). The dashed curves show $E_{T=0}$ and $E_{T=1}$ as functions of the pion coupling constant f_π^2 (the scale is given in the upper line of the box) for the ρ-meson constant f_ρ^2 taken in the "weak meson" limit. The collapse of the RPA frequency $E_{T=0}$ occurs when the actual value of the pion constant $f_\pi^2 \approx 0.08$ is used.

Fig. 2. a) Energies of the 0^- states $E_{T=0}$ and $E_{T=1}$ [Eq.(23)], in units of ω are plotted as a function of the ρ-meson strength f_ρ^2 calculated for the standard value of the π-meson constant. Collapse points in $E_{T=0}$ occur at values of f_ρ^2 approximately equal to the "strong" and "weak" ρ-meson couplings.

b) The enhancement factors F^P (crosses) and F^{PT} (solid lines) [Eq.(20)] as function of f_ρ^2 (for the standard value of f_π^2).